

The efficiency of the magnetic acceleration in relativistic jets

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Abstract. Using steady, axisymmetric, ideal magnetohydrodynamics (MHD) we analyze relativistic outflows by means of examining the momentum equation along the flow and in the transfield direction. We argue that the asymptotic Lorentz factor is $\gamma_\infty \sim \mu - \sigma_M$, and the asymptotic value of the Poynting-to-matter energy flux ratio – the so-called σ function – is given by $\sigma_\infty/(1 + \sigma_\infty) \sim \sigma_M/\mu$, where σ_M is the Michel’s magnetization parameter and μc^2 the total energy-to-mass flux ratio. We discuss how these values depend on the conditions near the origin of the flow. By employing self-similar solutions we verify the above result, and show that a Poynting-dominated flow near the source reaches equipartition between Poynting and matter energy fluxes, or even becomes matter-dominated, depending on the value of σ_M/μ .

Keywords: MHD, methods: analytical, relativity

1. Introduction

The main driving mechanism for relativistic outflows in AGNs, GRBs, and pulsar winds, is likely related to magnetic fields. These fields are able to tap the rotational energy of the disk, and accelerate matter ejecta not only magnetocentrifugally, but also due to the magnetic pressure. Using ideal magnetohydrodynamics (MHD) we examine initially Poynting-dominated outflows, trying to answer the following basic question: Which part of the total ejected energy flux is transferred to the matter kinetic energy flux asymptotically, and how this value depends on the conditions near the origin of the flow?

The system of equations of special relativistic, steady, cold, ideal MHD, consist of the Maxwell equations $0 = \nabla \cdot \mathbf{B} = \nabla \times \mathbf{E} = \nabla \times \mathbf{B} - 4\pi \mathbf{J}/c = \nabla \cdot \mathbf{E} - 4\pi J^0/c$, the Ohm’s law $\mathbf{E} = \mathbf{B} \times \mathbf{V}/c$, the continuity $\nabla \cdot (\rho_0 \gamma \mathbf{V}) = 0$, and momentum $-\gamma \rho_0 (\mathbf{V} \cdot \nabla) (\gamma \mathbf{V}) - \nabla P + (J^0 \mathbf{E} + \mathbf{J} \times \mathbf{B})/c = 0$ equations. Here \mathbf{V} is the velocity of the outflow, γ the associated Lorentz factor, (\mathbf{E}, \mathbf{B}) the electromagnetic field as measured in the central object’s frame, $J^0/c, \mathbf{J}$ the charge and current densities, and ρ_0 the gas rest-mass density in the comoving frame.

Assuming axisymmetry [$\partial/\partial\phi = 0$, in cylindrical (z, ϖ, ϕ) coordinates with \hat{z} along the rotation axis], four conserved quantities along the flow exist. If $A = (1/2\pi) \iint \mathbf{B}_p \cdot d\mathbf{S}$ is the poloidal magnetic flux func-



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tion, they are (e.g., Vlahakis & Königl, 2003a)¹: the mass-to-magnetic flux ratio $\Psi_A(A) = 4\pi\gamma\rho_0 V_p/B_p$, the field angular velocity $\Omega(A) = (V_\phi/\varpi) - (V_p/\varpi)(B_\phi/B_p)$, the specific angular momentum $L(A) = \gamma\varpi V_\phi - \varpi B_\phi/\Psi_A$, and the energy-to-mass flux ratio

$$\mu(A)c^2 = \gamma c^2 - \varpi\Omega B_\phi/\Psi_A. \quad (1)$$

The right-hand side of eq. (1) consists of the matter energy-to-mass flux ratio γc^2 , and the Poynting-to-mass flux ratio $(\mu - \gamma)c^2$. The sum of these two parts is a constant of motion, while their ratio is defined as the Poynting-to-matter energy flux ratio $\sigma = (\mu - \gamma)/\gamma$.

All the physical quantities can be written as functions of (A, σ) :

$$B_p = \frac{|\nabla A|}{\varpi}, \quad B_\phi = -\frac{\sigma}{1+\sigma} \frac{\mu\Psi_A c}{x}, \quad E = xB_p, \quad (2)$$

$$\gamma = \frac{\mu}{1+\sigma}, \quad \rho_0 = \frac{\sigma\Psi_A^2}{4\pi x^2} \left[1 - \frac{x_A^2}{x^2} + \sigma \frac{1-x_A^2}{x^2} \right]^{-1}, \quad (3)$$

$$\gamma \frac{V_p}{c} = \frac{\Omega^2 \varpi |\nabla A|}{\sigma \Psi_A c^3} \left[1 - \frac{x_A^2}{x^2} + \sigma \frac{1-x_A^2}{x^2} \right], \quad \gamma \frac{V_\phi}{c} = \mu \frac{x_A^2 - \frac{\sigma}{1+\sigma}}{x}, \quad (4)$$

where $x = \varpi\Omega/c$ is the cylindrical distance in units of the light cylinder's lever arm, and $x_A = (L\Omega/\mu c^2)^{1/2}$ its value at the Alfvén point.

The functions $A(\varpi, z)$, $\sigma(\varpi, z)$ obey the two remaining equations of the system: The Bernoulli equation² which is a quartic for σ

$$\left[\frac{\mu}{1+\sigma} \right]^2 = 1 + \left[\frac{\Omega^2 \varpi |\nabla A|}{\sigma \Psi_A c^3} \left(1 - \frac{x_A^2}{x^2} + \sigma \frac{1-x_A^2}{x^2} \right) \right]^2 + \left[\mu \frac{x_A^2 - \frac{\sigma}{1+\sigma}}{x} \right]^2, \quad (5)$$

and the transfield force-balance equation

$$\begin{aligned} & \left[1 - \frac{x_A^2}{x^2} \right] \left[1 + \frac{1}{\sigma} \right] \frac{\varpi \bar{L} A}{|\nabla A|} + 2 \frac{\hat{\varpi} \cdot \nabla A}{|\nabla A|} - \frac{\mu \Psi_A c^6}{\Omega^3 \varpi |\nabla A|} \left[\frac{\sigma}{1+\sigma} \right]^2 \frac{d}{dA} \frac{\mu \Psi_A}{\Omega} \\ & + \frac{\varpi |\nabla A|}{\Omega} \frac{d\Omega}{dA} - \frac{1}{\sigma} \left[1 - \frac{x_A^2}{x^2} + \sigma \frac{1-x_A^2}{x^2} \right] \frac{\varpi \nabla A \cdot \nabla \ln \frac{|\nabla A|}{\varpi}}{|\nabla A|} \\ & + \frac{\sigma \mu^2 c^4 \Psi_A^2 \left[x_A^2 - \frac{\sigma}{1+\sigma} \right]^2 \hat{\varpi} \cdot \nabla A}{x^4 \Omega^2 |\nabla A|^3 \left[1 - \frac{x_A^2}{x^2} + \sigma \frac{1-x_A^2}{x^2} \right]} - \frac{\mu^2 c^4 \Psi_A^2 \sigma \varpi \nabla A \cdot \nabla \sigma}{x^2 \Omega^2 |\nabla A|^3 (1+\sigma)^3} = 0. \quad (6) \end{aligned}$$

Here the operator $\bar{L} \equiv \nabla^2 - \frac{2}{\varpi} \hat{\varpi} \cdot \nabla$ is related to the curvature radius of the poloidal field lines $\mathcal{R} = |\nabla A| (\bar{L} A - \nabla A \cdot \nabla \ln |\nabla A/\varpi|)^{-1}$.

¹ The subscripts p/ϕ denote poloidal/azimuthal components.

² This equation comes from the identity $\gamma^2 = 1 + (\gamma V_p/c)^2 + (\gamma V_\phi/c)^2$.

In the force-free limit $\sigma = \infty$ it is $\mu = \infty$, $\Psi_A = 0$, $L = \infty$, $x_A = 1$, while μ/σ , $\mu\Psi_A$, $L\Psi_A$, $\sigma(1 - x_A^2)$ are finite. In this case, only the first four terms of eq. (6) survive, resulting in the “pulsar equation”

$$\left[1 - \frac{1}{x^2}\right] \frac{\varpi \bar{L} A}{|\nabla A|} + 2 \frac{\hat{\omega} \cdot \nabla A}{|\nabla A|} - \frac{\mu \Psi_A c^6}{\Omega^3 \varpi |\nabla A|} \frac{d}{dA} \frac{\mu \Psi_A}{\Omega} + \frac{\varpi |\nabla A|}{\Omega} \frac{d\Omega}{dA} = 0. \quad (7)$$

2. The σ function and its asymptotic value

2.1. A GENERAL ANALYSIS

An important combination of the field line constants is the “Michel’s magnetization parameter” $\sigma_M(A) = A\Omega^2/\Psi_A c^3$. In terms of σ_M , and using eqs. (1) and (2), we may write the exact expression³

$$\boxed{\frac{\sigma}{1 + \sigma} = \left(\frac{\sigma_M}{\mu}\right) \left(\frac{-B_\phi}{E}\right) \left(\frac{B_p \varpi^2}{A}\right)} \quad (8)$$

The left-hand side represents the Poynting-to-total energy flux, and – using the first of eqs. (3) – can be rewritten as $(\mu - \gamma)/\mu$. As long as the flow is Poynting-dominated ($\gamma \ll \mu$), this ratio is close to unity, and $\sigma \gg 1$. This continues to be the case in the neighborhood of the classical fast magnetosound surface, where $\gamma \approx \mu^{1/3}$ (e.g., Camenzind, 1986). As a result, the superfast regime of the flow is the only place where a transition from high ($\gg 1$) to ~ 1 or $\ll 1$ values of σ is possible. In this regime, and for extremely relativistic flows, the term $(-B_\phi/E)$ is very close to unity.⁴ Hence, eq. (8) gives a simple relation between the σ function and the ratio $B_p \varpi^2/A$:

$$\frac{\sigma}{1 + \sigma} \approx \left(\frac{\sigma_M}{\mu}\right) \left(\frac{B_p \varpi^2}{A}\right). \quad (9)$$

Suppose that the value of the function $B_p \varpi^2/A$ near the classical fast surface is $(B_p \varpi^2/A)_f$. Since $\sigma \gg 1$ at this point, eq. (9) implies that the constant of motion $\sigma_M/\mu \approx 1/(B_p \varpi^2/A)_f$.

Denoting with δl_\perp the distance between two neighboring poloidal field lines A and $A + \delta A$, magnetic flux conservation implies $B_p \varpi^2 = (\varpi/\delta l_\perp) \delta A$. Thus, a decreasing $B_p \varpi^2$ – and hence, from eq. (9), a

³ Eq. (8) remains the same with thermal effects included.

⁴ The requirement that the Lorentz invariant $B^2 - E^2 > 0$, using $E = x B_p$, gives $B_\phi^2/E^2 > 1 - 1/x^2$. In addition, eqs. (3)–(4) give $V_\phi/c = x + (V_p/c)(B_\phi/B_p)$, and the condition $V_\phi > 0$ implies $-B_\phi/E < c/V_p$. Thus, $(1 - 1/x^2)^{1/2} < -B_\phi/E < c/V_p$.

decreasing σ function – corresponds to poloidal field lines expanding in a way such that their distance δl_\perp increases faster than ϖ . How fast the field lines expand is determined by the transfield force balance equation; thus, eq. (6) indirectly determines the flow acceleration. Since the available solid angle for expansion is finite, there is a minimum value of the $B_p \varpi^2/A$ function. The field lines asymptotically have a shape $z \approx z_0(A) + \varpi/\tan \vartheta(A)$, where $\vartheta(A)$ is their opening angle. Differentiating the latter equation we get $B_p \varpi^2/A = (A\vartheta'/\sin \vartheta - Az'_0 \sin \vartheta/\varpi)^{-1}$, a decreasing function, reaching a minimum value $\sin \vartheta/A\vartheta'$ at $\varpi \gg z'_0 \sin^2 \vartheta/\vartheta'$ (Vlahakis, 2004). Since the factor $\sin \vartheta/A\vartheta'$ is ~ 1 , the minimum value of the $B_p \varpi^2/A$ function is ~ 1 , corresponding to⁵

$$\frac{\sigma_\infty}{1 + \sigma_\infty} \approx \frac{\sigma_M}{\mu} \left(\frac{B_p \varpi^2}{A} \right)_\infty \approx \frac{(B_p \varpi^2)_\infty}{(B_p \varpi^2)_f} \sim \frac{\sigma_M}{\mu} \sim \frac{1}{(B_p \varpi^2/A)_f}. \quad (10)$$

Equivalently, the asymptotic Lorentz factor is $\gamma_\infty = \mu/(1 + \sigma_\infty) \sim \mu - \sigma_M$, and the asymptotic Poynting-to-mass flux ratio is $\sim \sigma_M c^2$.

Another interesting connection with the boundary conditions near the source can be found, by noting that, as long as $|B_\phi| \approx E = xB_p$, $B_p \varpi^2/A \approx 2|I|/A\Omega$, where $|I| = (c/2)\varpi|B_\phi|$ is the poloidal current. Thus, $(B_p \varpi^2/A)_f \approx 2|I|_f/A\Omega$, and since $|I|$ remains constant of motion inside the force-free subfast regime, $(B_p \varpi^2/A)_f \approx 2|I|_i/A\Omega$, and $\mu/\sigma_M \approx 2|I|_i/A\Omega$. Hence, eq. (10) implies a direct connection of the efficiency and the Lorentz factor to the ejection characteristics

$$\frac{\sigma_\infty}{1 + \sigma_\infty} \approx \frac{A\Omega}{2|I|_i} \left(\frac{B_p \varpi^2}{A} \right)_\infty \sim \frac{A\Omega}{2|I|_i}, \quad \gamma_\infty \sim \mu \left(1 - \frac{A\Omega}{2|I|_i} \right) \quad (11)$$

2.2. THE VALUE OF σ_∞ IN KNOWN SOLUTIONS

Solving the system of equations (5) and (6) is highly intractable. In the following we review the currently known methods to obtain solutions.

2.2.1. Numerical methods

Eq. (5) is a relatively simple algebraic equation for σ . After substituting its solution (in terms of A and its derivatives) in the transfield

⁵ The only exception to this general result is to have asymptotically $\sigma/(1 + \sigma) \ll \sigma_M/\mu$ in some finite solid angle regions, combined with other regions with bunched field lines [in which $B_p \varpi^2 \gg A$ and $\sigma/(1 + \sigma) \gg \sigma_M/\mu$]. Note also that the most general asymptotic field line shape slightly deviates from straight lines, resulting in a logarithmic acceleration reaching σ values smaller than σ_M/μ (Chiueh, Li, & Begelman, 1991; Okamoto, 2002; Vlahakis, 2004). However, this acceleration can happen in exponentially large distances and hence is physically irrelevant.

force-balance equation (6), we get a second order partial differential equation for the magnetic flux function A . Its solution determines the field-streamline shape on the poloidal plane. Due to the fact that this equation is of mixed type, i.e., changes from elliptic to hyperbolic, it is beyond the capability of existing numerical codes to solve this highly nonlinear problem, and no solution has been obtained so far.

An alternative numerical approach is to solve the time-dependent problem (hyperbolic in time) and expect to reach a steady-state. However, all existing codes fail to simulate relativistic magnetohydrodynamic flows for more than a few rotational periods. On top of that, it is not always clear how the issue of the boundary conditions is handled.

A promising combination of the two above methods is followed by Bogovalov (2001), who solves the inner problem using time-dependent evolution (avoiding the elliptic to hyperbolic transitions), and the outer problem using steady-state equations. The code is not yet capable of solving the problem at large distances, though.

2.2.2. *The force-free assumption*

In the force-free limit $\sigma = \infty$, the two equations (5) and (6) decouple. Thus, one may solve the somewhat simpler, elliptic equation (7) (e.g., Contopoulos et al., 1999), and then solve eq. (5) for σ .

The force-free solutions have several problems, the most important of which is the following: Since the back reaction of the matter to the field is neglected, the drift velocity soon after the light cylinder becomes larger than the light speed. Equivalently, there is no frame of reference where the electric field vanishes, and eq. (5) has no real solutions for σ .

In any case, the force-free assumption brakes down in the superfast ($\gamma > \mu^{1/3} \Leftrightarrow \sigma < \mu^{2/3} - 1$) regime, where the flow becomes hyperbolic and the back reaction of the matter to the field cannot be neglected. Since for Poynting-dominated flows the value of the σ function at the classical fast surface is $\mu^{2/3} - 1 \gg 1$, the force-free method cannot be used for examining the efficiency of the magnetic acceleration.

2.2.3. *The prescribed field line shape assumption*

If one assumes a known magnetic flux distribution, i.e., a known function $B_p \varpi^2 / A = \varpi |\nabla A| / A$, then it is trivial to solve eq. (5) for σ and find the flow speed⁶ (e.g., Takahashi et al., 1990; Fendt & Greiner, 2001). Thus, when we use this method, practically we implicitly give the function σ ! However, these solutions do not satisfy the transfield force-balance equation (6); thus, they are not fully self-consistent.

⁶ In the superfast regime, this equation reduces to the much simpler eq. [9].

2.2.4. *The monopole approximation*

This is a subcase of the prescribed field line case, based on the assumption that the poloidal magnetic field is quasi-monopolar. This assumption, with the help of eq. (9), is equivalent to the assumption that the σ function is constant! In fact, a tiny acceleration is possible in the subfast regime, leading to $\gamma_\infty \sim \mu^{1/3}$ and $\sigma_\infty \sim \mu^{2/3} \gg 1$ (Michel, 1969). This solution gave the erroneous impression to the community that relativistic MHD is in general unable to give high acceleration efficiencies. However, the solution corresponds to a special case of boundary conditions, and most importantly, it does not satisfy the transfield force-balance equation.

2.2.5. *Other approximate solutions*

Various tries to solve a simplified version of eqs. (5) and (6), by neglecting some terms [e.g., Tomimatsu & Takahashi (2003) who neglected the curvature radius term in eq. (6)] involve a risk, since the system of equations is highly nonlinear and it could be crucial to keep terms that seem at first negligible, especially second order terms.

Also perturbations around a monopolar solution [e.g., Beskin et al. (1998) who perturbed a solution with $\gamma = \infty$ and $V_\phi = 0$] are (using eq. [9]) equivalent to perturbations of σ around a constant high value.

Concluding, **in order to solve the efficiency problem, one has to solve simultaneously eqs. (5) and (6).**

2.2.6. *The r self-similar special relativistic model*

The only known exact solution of eqs. (5) and (6) is the r self-similar special relativistic model, found independently by Li, Chiueh, & Begelman (1992) and Contopoulos (1994) in the cold limit, and further generalized by Vlahakis & Königl (2003a) including thermal effects. It corresponds to boundary conditions in a conical surface ($\theta = \theta_i$ in spherical coordinates $[r, \theta, \phi]$) of the form $B_r = C_1 r^{F-2}$, $B_\phi = -C_2 r^{F-2}$, $V_r = C_3$, $V_\theta = -C_4$, $V_\phi = C_5$, $\rho_0 = C_6 r^{2(F-2)}$, $P = C_7 r^{2(F-2)}$, with constant C_1, \dots, C_7 . The parameter of the model F controls the initial current distribution [$-\varpi B_\phi = C_2 \sin \theta_i r^{F-1}$ is an increasing or decreasing function of r for $F >$ or < 1 ; see Vlahakis & Königl (2003a) for details]. Despite the assumed form of the boundary conditions, the assumption that gravity is negligible, and the absence of intrinsic scale, r self-similar remain the only self-consistent relativistic MHD solutions.

Figure 1 shows the function $B_p \varpi^2 / A$ for various r self-similar solutions with application to GRB and AGN outflows. With the help of the figure and the known μ , μ/σ_M values, we are able to fill the table

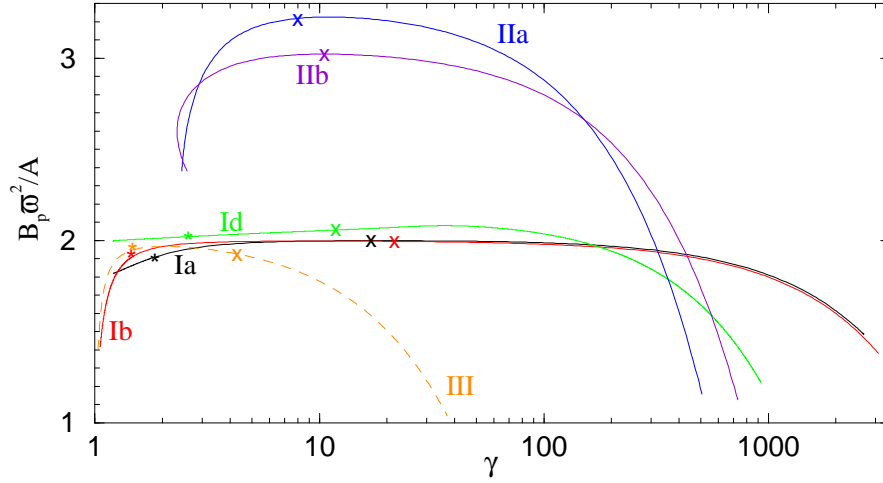


Figure 1. Function $B_p \varpi^2 / A$ versus γ for exact r self-similar solutions presented in Vlahakis & Königl (2003a, 2003b, 2004). The positions of the Alfvén and classical fast magnetosound points are marked with “*” and “x”, respectively.

Table I. The efficiency in r self-similar models

Solution	μ	$\frac{\mu}{\sigma_M}$	$\left(\frac{B_p \varpi^2}{A}\right)_f$	$\left(\frac{B_p \varpi^2}{A}\right)_\infty$	$\frac{(B_p \varpi^2)_\infty}{(B_p \varpi^2)_f}$	$\frac{\sigma_\infty}{1 + \sigma_\infty}$
Ia ^a	10116.1	2.02	1.99	1.49	0.74	0.73
Ib ^a	9997.4	2.00	1.99	1.38	0.69	0.69
Id ^a	2150.0	2.15	2.06	1.22	0.59	0.57
IIa ^b	778.9	3.31	3.22	1.16	0.36	0.35
IIb ^b	1156.6	3.06	3.02	1.13	0.37	0.37
III ^c	75.0	2.06	1.93	1.04	0.54	0.51

^a Trans-Alfvénic GRB outflow solutions, from Vlahakis & Königl (2003a).

^b Super-Alfvénic GRB outflow solutions, from Vlahakis & Königl (2003b).

^c AGN jet solution, from Vlahakis & Königl (2004).

I.⁷ This table verifies the results of § 2, and in particular: 1) Since the values of the 3rd and 4th columns are approximately equal, $\sigma_M/\mu \approx 1/(B_p \varpi^2/A)_f$ indeed holds. 2) The value $(B_p \varpi^2/A)_\infty$ is indeed of order unity. 3) From the two last columns we verify that the asymptotic value of σ satisfies eq. (9), and $\sigma_\infty/(1 + \sigma_\infty) \approx (B_p \varpi^2/A)_f^{-1} (B_p \varpi^2/A)_\infty \approx (\sigma_M/\mu)(B_p \varpi^2/A)_\infty$. 4) The value $\sigma_\infty/(1 + \sigma_\infty)$ is roughly equal to σ_M/μ . 5) The value γ_∞ is roughly equal to $\mu - \sigma_M = \mu(1 - \sigma_M/\mu)$. A common characteristic of all trans-Alfvénic r self-similar solutions is that $\mu/\sigma_M \sim 2$, and thus $\gamma_\infty \sim \mu/2$.

⁷ Besides the values of $B_p \varpi^2/A$, fig. 1 shows γ_∞ ; we also find $\sigma_\infty = (\mu/\gamma_\infty) - 1$.

2.2.7. z self-similar asymptotic solutions

Vlahakis (2004) derived z self-similar asymptotic solutions of the system of eqs. (5) and (6) that verify the presented analysis. For conditions near the classical fast magnetosound surface corresponding to $(B_p \varpi^2)_f \gg A \Leftrightarrow \sigma_M \ll \mu \Leftrightarrow |I|_i \gg A\Omega/2$ the σ_∞ is found to be $\ll 1$.

3. Conclusion

- In order to solve for the acceleration it is absolutely necessary to solve for the poloidal field line shape as well. The Bernoulli and trans-field force-balance equations are interrelated and we cannot solve them separately, especially in the superfast regime.
- Models that assume quasi-monopolar magnetic field (the Michel's solution included), equivalently *assume* that the magnetic acceleration is inefficient. Moreover, no fully consistent solution of this kind is known.
- $\sigma_\infty/(1 + \sigma_\infty) \approx (\sigma_M/\mu)(B_p \varpi^2/A)_\infty \sim \sigma_M/\mu$ is an analytic expression for the asymptotic Poynting-to-total energy flux ratio, confirmed by self-similar solutions. Possible acceleration to smaller σ_∞ can only happen in exponentially large (and thus physically irrelevant) distances.
- $\sigma_\infty/(1 + \sigma_\infty) \approx (A\Omega/2|I|_i)(B_p \varpi^2/A)_\infty \sim A\Omega/2|I|_i$ is a general result connecting the asymptotic value σ_∞ (and hence $\gamma_\infty = \mu/[1 + \sigma_\infty]$) with the conditions near the origin of an initially Poynting-dominated flow.
- The asymptotic Lorentz factor is $\gamma_\infty \sim \mu - \sigma_M \approx \mu(1 - A\Omega/2|I|_i)$, and the asymptotic Poynting-to-mass flux ratio is $\sim \sigma_M c^2 \approx \mu c^2 A\Omega/2|I|_i$.

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